Nuclear Coherent versus Incoherent Effects in Peripheral RHI Collisions.

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Abstract

We derive simple and physically transparent expressions for the contribution of the strong interaction to one nucleon removal processes in peripheral relativistic heavy ion collisions. The coherent contribution,i.e, the excitation of a giant dipole resonance via meson exchange is shown to be negligible as well as interference between coulomb and nuclear excitation. Incoherent nucleon knock out contribution is also derived suggesting the nature of the nuclear interaction in this class of processes. We also justify the simple formulae used to fit the data of the E814 Collaboration.

1   Introduction

The important role played by the electromagnetic interaction in single nucleon removal processes induced by Relativistic Heavy Ions (RHI) has long been known. Dissociation of RHI by the Coulomb field of the target nucleus was first reported by Heckman and Lindstron [1] using for that purpose radiochemical techniques. Many other accelerator experiments have followed, using several combinations of heavy projectiles and targets with incoming energies $E \gtrsim 1$. GeV/nucleon. The most conspicuous evidence of such electromagnetic processes is their large cross section and approximate dependence with

Z^2 , Z being the atomic number. Within this context it is very important to have a quantitative model which allows one to isolate and identify possible contributions arising from different processes due, for instance, to the strong interaction.

Several simple models have been used to separate the nuclear contribution the most popular among experimental groups being the Limiting Fragmentation Model (LFM). This model assumes that the yield of the particular fragment from the target due to nuclear interactions will be independent of the beam, except for a geometrical factor [2]-[7]. Based on this model a parametrization of the nuclear contribution is performed and the electromagnetic contribution extracted from the data. More recently in Brookhaven the same type of process is investigated [8],[9]. Their data analysis assumes the one nucleon removal to be basically due to Coulomb dissociation (parametrized as bZ_T^c, b and c being free parameters) and to nuclear effects (parametrized as the sum of the heavy ions' radii $(a \ (A_T^{1/3} + A_P^{1/3})$, a being a free parameter). In this context an essential question arises: is it possible to give a sound theoretical basis for the contribution of the strong interaction in the one nucleon removal process? Can one theoretically justify the E814 prescription, for the nuclear cross section, quantitatively? Answering these questions are the main purpose of the present work.

We consider two mechanisms which could be responsible for one nucleon removal me-diated by the nuclear field: the excitation of a giant dipole resonance via vector meson exchange and the direct incoherent knock-out contribution. In the first case, i.e., for the collective effects we use the framework of ref.[10] and derive analytical expressions for pure Coulomb (rather well known), pure nuclear and interference cross sections (here obtained for the first time). We show that the theoretical cross sections involving the strong inter-action are independent of the bombarding energy, depend very sensitively on the meson mass and display a different qualitative behavior as compared to the empirically predicted one from E814. Moreover for realistic values of the coupling constant and mesonic masses these cross sections are too small. We are therefore led to the conclusion that one nucleon removal through the excitation of a giant dipole resonance by means of the nuclear field cannot be the mechanism responsible for the observed nuclear contribution.

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We next investigate the possibility that one nucleon is removed by means of and incoherent process (fragmentation reaction). The adequate theoretical tool for this purpose is Glauber's multiple scattering theory [11]. It has long been known both for its simplicity and amazing predictive power. One can find copious examples in the literature where such theory allows for a simple physical interpretation of experimental results as well as their quantitative analysis ([12]-[14]). We therefore work in such scheme. In fact, fragmentation reactions of the type discussed here have already been successfully analyzed in the framework of Glauber's theory: in one nuclear removal reactions, the momentum distribution of the outgoing fragment has been shown to reflect the momentum distribution of the nucleon which is removed from the surface of the projectile nucleus [14]. This is precisely the physics behind the incoherent nuclear one nucleon removal, as we shall show here. The target dependence so obtained corresponds precisely to the empirically fitted one. We are thus able to reproduce the results of the E814 experiment and to give a physical interpretation of the empirical fit. The calculation presented here, we believe, provides for a good theoretical prescription as to the quantitative nuclear contribution to the process.

This paper is organized as follows: in section 2 the Coherent Effects, i.e, the excitation of a Giant Dipole Resonance via one photon and one vector meson exchange are analyzed. Analytical expressions for the pure nuclear and interference contributions are derived. Section 3, presents the study of the Nuclear Incoherent Effects, using Glauber's formalism. Finally we compare our results with E814 data showing very good agreement.

2   Nuclear Coherent Effects

Recently the E814 Collaboration [8] has reported on the total cross section for the excitations of the giant dipole resonance by means of peripheral heavy ion collisions. The reaction considered was ^{28}Si into $p +^{27} Al$ at an energy of 14.5 GeV/nucleon bombard-

ing different targets (C, Al, Cu, Sn, Pb). The experimental data suggest a non negligible amount of nuclear effects and show that the final-state energy, is peaked near the isovector giant dipole resonance in ^{28}Si.

The pure Coulomb contribution is quantitatively the most important and has been previously obtained by several authors by different methods (Refs. [10], [15],[16]). On the other hand the pure nuclear and interference cross sections are not as well understood (see ref. [17]).

In this section we study the excitation of a giant dipole resonance via one photon and one meson exchange and assume the one nucleon removal to be entirely due to this process.

The study of vector and scalar meson exchanges in a covariant formalism has been performed in detail in ref. [18] for inelastic proton-nucleus scattering. In this reference it is shown that within the relativistic eikonal approximation one can sum an infinite series of Feynman diagrams and the resulting scattering amplitude is strongly dominated by the vertex where the excitation takes place. We therefore restrict ourselves in our derivations to the lowest order Feynman diagram contributions for the excitation of a giant dipole resonance via vector meson and photon exchange.

We show that in the ultra-relativistic limit it is possible to derive very simple analytical formula for the three contributions (pure Coulomb, pure nuclear and interference), to understand the physical reason why the pure nuclear contribution is so much smaller than the Coulomb one and to estimate the contribution of interference.

2.1   Collective Coulomb plus Nuclear Excitation Cross Section

Formally, the first order contribution to collective nuclear excitation via one photon and one meson exchange between projectile and target nucleus is given in the Center-of-Mass System (C.M.) by [10]:

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$$\left.\frac{d\sigma}{d\Omega}\right|_{CM} = \sum_{f_A}\sum_{f_B}\frac{M'_A M'_B M_A M_B}{(\sqrt{s})^2}\left[\frac{2\alpha Z_A Z_B}{q^2}+\frac{g_v}{q^2-\mu^2}\right]^2$$
$$|F_B^\mu(\vec{q}_B,\lambda_B,m_B)\Lambda_\nu^\mu(P_A\to P_B)F_A^\nu(\vec{q}_A,\lambda_A,m_A)|^2 \tag{1}$$

where α is the fine-structure-constant, $\sqrt{s}$ is the total center of mass energy, the Z's and M's are atomic and mass numbers of each nucleus, respectively, q is the four-momentum transfer and $\vec{q}_A$ ($\vec{q}_B$) its spatial part in the proper system of the nucleus A (B) (ref. [10]). g_v is the strong coupling constant and μ is the vector meson's mass.

The F_A^l are the Fourier transforms of transition current matrix elements, given by:

$$\begin{aligned}
F_A^l(q_A;\lambda_A,m_A) &= \frac{1}{C}\int d^3r <E_A^*;\lambda_A,m_A|j_A^l(\vec{r})|Ag.s> e^{i\vec{q}_A\cdot\vec{r}}\\
&= <E_A^*;\lambda_A,m_A|j_A^l(q_A)|A_{g.s}> \tag{2}
\end{aligned}$$

with $C=eZ_A$ for the Coulomb contribution and $C=g_v$ for the strong contribution. We can write the similar expression for the nucleus B.

In the above equations, currents and nuclear states are represented in their proper systems, and $|E;\lambda,m>$ stands for the final state of nucleus A or B with angular momentum λ whose projection in the incident beam direction (e'_z) is m. $\sum_f$ means the summation over final states. The Lorentz transformation matrix is given by

$$\Delta_{\tau l} = \begin{pmatrix} \gamma & 0 & 0 & -\beta\gamma \\ 0 & -1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ \beta\gamma & 0 & 0 & -\gamma \end{pmatrix} \tag{3}$$

where β is the velocity of nucleus A in the rest frame B, and γ is the associated Lorentz factor ($\gamma=\sqrt{1-\beta^2}$).

In order to evaluate the current matrix elements, it is convenient to use the coordinate system whose z–direction is taken to be the direction of the momentum transfer seen by the nucleus. The cartesian base vectors $\{\hat{e}_x,\hat{e}_y,\hat{e}_z\}$, where $\hat{z}=q/|\vec{q}|$, are related to the original ones $\{\hat{e}'_x,\hat{e}'_y,\hat{e}'_z\}$ by

$$
\begin{pmatrix} \vec{e_x} \\ \vec{e_y} \\ \vec{e_z} \end{pmatrix} = \begin{pmatrix} (\cos\theta - 1)\cos^2\varphi + 1 & (\cos\theta - 1)\sin\varphi\cos\theta & -\sin\theta\cos\varphi \\ (\cos\theta - 1)\sin\varphi\cos\varphi & (\cos\theta - 1)\sin^2\varphi + 1 & -\sin\theta\sin\varphi \\ \sin\theta\cos\varphi & \sin\theta\sin\varphi & \cos\theta \end{pmatrix}
$$
$$
\times \begin{pmatrix} \vec{e'_x} \\ \vec{e'_y} \\ \vec{e'_z} \end{pmatrix} \tag{4}
$$

where θ and φ are polar angles of $\hat{z}$ with respect to the base vectors $\{\hat{x}', \hat{y}', \hat{z}'\}$. The corresponding spherical base vectors $\{\hat{e}_{-1}, \hat{e}_0, \hat{e}_{+1}\}$ and $\{\hat{e}'_{+1}, \hat{e}'_0, \hat{e}'_{-1}\}$ are related to each other by

$$
\begin{pmatrix} \hat{e}_{+1} \\ \hat{e}_0 \\ \hat{e}_{-1} \end{pmatrix} = \begin{pmatrix} \hat{e}'_{+1} \\ \hat{e}'_0 \\ \hat{e}'_{-1} \end{pmatrix} D^{(1)}(\varphi, \theta, -\varphi) \tag{5}
$$

where $D^{(1)}(\varphi, \theta, -\varphi)$ is the usual rotation matrix for total angular momentum $j = 1$. Using these relations we can express the four-vector current in terms of spherical components as

$$
\begin{pmatrix} \rho \\ j'_x \\ j'_y \\ j'_z \end{pmatrix} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & & & \\ 0 & & UD^T & \\ 0 & & & \end{pmatrix} \begin{pmatrix} \rho \\ j_{+1} \\ j_0 \\ j_{-1} \end{pmatrix} \tag{6}
$$

where the 3×3 matrix U transforms spherical vectors into cartesian ones. The convenience of the spherical basis now becomes apparent: the selection rules for the current matrix elements are

$$
< E^*; \lambda, \nu | \rho | 0, 0 > = < \rho > \delta_{\nu,0},
$$

$$
< E^*; \lambda, \nu | j_{+1} | 0, 0 > = < j_{+1} > \delta_{\nu,+1},
$$

$$
< E^*; \lambda, \nu | j_0 | 0, 0 > = < j_0 > \delta_{\nu,0}, \tag{7}
$$

$$
< E^*; \lambda, \nu | j_{-1} | 0, 0 > = < j_{-1} > \delta_{\nu,-1} = < j_{+1} > \delta_{\nu,-1},
$$

where ν refers to the projection of angular momentum in the three $\vec{q}$ directions.

We can now write the matrix element in eq. (1) as

$$
M = \bar{D}^{(\lambda_B)}(\hat{q}_B)
\begin{pmatrix}
0 & < j_{+1}^B > & 0 & 0 \\
< \rho^B > & 0 & < j_0^B > & 0 \\
0 & 0 & 0 & < j_{-1}^B >
\end{pmatrix} \times
$$

$$
\begin{pmatrix}
1 & 0 & 0 & 0 \\
0 & & & \\
0 & & D^{(1)\dagger}(\hat{q}_B) & \\
0 & & &
\end{pmatrix} \times
\begin{pmatrix}
\gamma & 0 & -\beta\gamma & 0 \\
0 & 0 & 0 & 1 \\
\beta\gamma & 0 & -\gamma & 0 \\
0 & 1 & 0 & 0
\end{pmatrix} \times
\begin{pmatrix}
1 & 0 & 0 & 0 \\
0 & & & \\
0 & & D^{(1)*}(\hat{q}_A) & \\
0 & & &
\end{pmatrix}
$$

$$
\times
\begin{pmatrix}
0 & < \rho^A > & 0 \\
< j_{+1}^A > & 0 & 0 \\
0 & < j_0^A > & 0 \\
0 & 0 & < j_{-1}^A >
\end{pmatrix}
\bar{D}^{(\lambda_A)T}(\hat{q}_A)
\tag{8}
$$

$\hat{q}$ (the argument of $D(\hat{q})$), symbolically represents the Euler angles associated to the rotation required to transform $\hat{e}'_z$ to $\hat{e}_z$ in each nucleus.

In the case where the nucleus B stays unexcited, we obtain

$$
\left.\frac{d\sigma}{d\Omega}\right|_{m=\pm1}^{single} = \frac{1}{2}F_c| < \rho_B > |^2\gamma^2
$$

$$
\sin^2\theta_A|[< \rho_A > -\beta\cos\theta_A < j_0^A >] + \beta\cos\theta_A < j_1^A > |^2
\tag{9}
$$

$$
\left.\frac{d\sigma}{d\Omega}\right|_{m=0}^{single} = F_c| < \rho_B > |^2\gamma^2
$$

$$
|\cos\theta_A\{< \rho_A > -\beta\cos\theta_A < j_0^A >\} - \frac{1}{2}\beta\sin^2\theta_A < j_1^A > |^2
\tag{10}
$$

where θ_A represents the polar angle of $\vec{q_A}$ in the rest frame of A and similar expressions can be constructed for the strong and interference part with the respective substitution of F_c by F_s and F_{cs}.

The next problem one has to face in the present formulation is the inclusion of nuclear absorption effects. The way to circumvent this problem is the following: we first write the total cross-section in terms of an integral over the transverse momentum transfer,

$$
\sigma_{tot} = \int d^2q_\perp \left[\frac{\sqrt{F_c}}{q^2} + \frac{\sqrt{F_s}}{q^2 - \mu^2}\right]^2 Tr[M^\dagger(\vec{q}_\perp)M(\vec{q}_\perp)]
\tag{11}
$$

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where $\sqrt{F_c} = 2\alpha Z_A Z_B$ and $\sqrt{F_s} = g_v$. We then introduce a Fourier transform

$$f(b) = \frac{1}{2\pi} \int d^2\vec{q}_\perp \left[\frac{\sqrt{F_c}}{q^2} + \frac{\sqrt{F_f}}{q^2 - \mu^2} \right] \sqrt{M(\vec{q}_\perp)} e^{i\vec{q}_\perp \cdot \vec{b}} \tag{12}$$

where $M(\vec{q}_\perp)$ denotes

$$
\begin{aligned}
M(\vec{q}_\perp) &= \sqrt{\frac{d^2\Omega}{dq_\perp^2}} M(\vec{q}) \\
&= \frac{1}{\sqrt{p\left(\sqrt{p^2 - q_\perp^2}\right)}} M(\vec{q}) \tag{13}
\end{aligned}
$$

and p stands for the projectile's center-of-mass momentum. The total cross section can now be written as

$$\sigma_{tot} = \int d^2 b\, Tr[f^\dagger(b) f(b)] \tag{14}$$

The above formula can be interpreted as the impact parameter representation of the cross section. It is now possible to estimate nuclear absorption effects in a simple way: we attribute those parts of the integral coming from all values of b smaller than the sum of radii of projectile and target to the absorption coming from nuclear interactions which do not contribute to excite the collective mode. Thus, with $b_{min} = r_0(A_T^{1/3} + A_P^{1/3})$, we get:

$$
\begin{aligned}
\sigma_{tot} &= \frac{1}{(2\pi)^4} \int_{b_{min}}^{\infty} d^2 b \int d^2 q_\perp \int d^2 q'_\perp \left[\frac{F_c}{q^2(q_\perp) q^2(q'_\perp)} + \right. \\
&\quad \left. 2\frac{\sqrt{F_c F_s}}{q^2(q_\perp)(q^2(q'_\perp) - \mu^2)} + \frac{F_s}{(q^2(q_\perp) - \mu^2)(q^2(q'_\perp) - \mu^2)} \right] \\
&\quad Tr\left[M^\dagger(\vec{q}_\perp) e^{-i(\vec{q}_\perp - \vec{q}'_\perp)\cdot\vec{b}} M(\vec{q}'_\perp) \right] \tag{15}
\end{aligned}
$$

where

$$
\begin{aligned}
\sigma_c &= \frac{1}{(2\pi)^4} \int_{b_{min}}^{\infty} d^2 b \int d^2 q_\perp \int d^2 q'_\perp \left[\frac{F_c}{q^2(q_\perp) q^2(q'_\perp)} \right] \\
&\quad Tr\left[M^\dagger(\vec{q}_\perp) e^{-i(\vec{q}_\perp - \vec{q}'_\perp)\cdot\vec{b}} M(\vec{q}'_\perp) \right], \tag{16}
\end{aligned}
$$

8

$$\sigma_s = \frac{1}{(2\pi)^4} \int_{b_{min}}^{\infty} d^2b \int d^2q_\perp \int d^2q'_\perp \left[\frac{F_s}{(q^2(q_\perp) - \mu^2)(q^2(q'_\perp) - \mu^2)} \right]$$
$$Tr\left[M^\dagger(\vec{q}_\perp) e^{-i(\vec{q}_\perp - \vec{q}'_\perp)\cdot\vec{b}} M(\vec{q}'_\perp) \right] \tag{17}$$

and

$$\sigma_{cs} = \frac{1}{(2\pi)^4} \int_{b_{min}}^{\infty} d^2b \int d^2q_\perp \int d^2q'_\perp \left[2\frac{\sqrt{F_c F_s}}{q^2(q_\perp)(q^2(q'_\perp) - \mu^2)} \right]$$
$$Tr\left[M^\dagger(\vec{q}_\perp) e^{-i(\vec{q}_\perp - \vec{q}'_\perp)\cdot\vec{b}} M(\vec{q}'_\perp) \right], \tag{18}$$

correspond to the Coulomb excitation process, to the pure strong contribution and to the interference between two, respectively.

2.2 Results

In order to obtain analytical results we now need a model for the matrix elements corresponding to the excitation of a giant dipole resonance. We use a macroscopic model. T. Suzuki and D. J. Rowe [19] have derive a sum rule for current matrix elements which corresponds to the incompressible fluid model. In it, the giant resonance state is represented by a unique level of given multipolarity and the transition matrix elements of the current from the ground state to this level is assumed to be of the form

$$< g.s.|J(\vec{r})|E^*, \lambda\nu >= \frac{N}{2im}\rho_0(\vec{r}) \bigtriangledown \{r^\lambda Y_{\lambda\nu}\} \tag{19}$$

where $\rho_0(\vec{r})$ is the ground state density distribution, m the nucleon mass and N the normalization constant related to the sum rule value S as,

$$N = \sqrt{E^*/S} \tag{20}$$

with

$$S = 3\lambda A R^{2\lambda-2}/8\pi m \qquad (21)$$

The Fourier transform of eq. (19) gives the required matrix element. After some algebra we obtain

$$< \rho >= \frac{N\sqrt{4\pi(2\lambda+1)}\lambda \int dr r^{\lambda+1}\rho_0(r)j_{\lambda-1}(qr)}{2mi^\lambda E^*/|\vec{q}|} \qquad (22)$$

where $j_{\lambda+1}(qr)$ are spherical Bessel functions and

$$< j_0 >= \frac{E^*}{|q|} < \rho >, \qquad (23)$$

$$< j_1 >= \left[\frac{\lambda+1}{2\lambda}\right]^{1/2} < j_0 > \qquad (24)$$

In this model, magnetic current is neglected. It is known that for giant resonances, the contribution of magnetic current is small. For very high energies ($\gamma >> 1$) and in the low momentum transfer limit ($qR << 1$), we may approximate the nuclear matrix element in eq. (22) as

$$< \rho >= \sqrt{\frac{AE^*}{2m}}\frac{|q_A|}{E^*} \qquad (25)$$

Furthermore, the four-momentum transfer squared is given by

$$q^2 \simeq -\left[\left(\frac{E^*}{\beta\gamma}\right)^2 + q_\perp^2\right] \qquad (26)$$

In this limit,

$$M_0(q_\perp) \sim 0 \qquad (27)$$

and

$$M_1(q_\perp) \sim \frac{\gamma}{p}(AE^*/2m)^{1/2}q_\perp e^{-iq} \qquad (28)$$

and it is possible to carry out the integrals in $q_\perp$ and $q'_\perp$ as well as over the impact parameter b. We obtain for the Coulomb contribution

$$\sigma_c = 3.5 \times 10^{-3} Z_B^2 \frac{N_A Z_A}{A_A^{2/3}} \ln\left[\delta/\xi + 1\right] mb. \tag{29}$$

where $\xi = E^* b_{min}/\beta\gamma$ and $\delta = 0.681085\ldots$.

The pure strong interaction and interference contributions can be obtained in an analogous manner. We get

$$\sigma_s = 26.65(2\pi^2)g_v^4 A_A^{4/3} e^{-2\mu b_{min}} \ mb \tag{30}$$

and

$$\sigma_{cs} = \frac{\alpha g_v^2 (\pi N_A Z_A)^{1/2}}{10m} Z_B A_B^{1/3} \frac{e^{-\mu b_{min}}}{\sqrt{\mu b_{min}}} \ mb \tag{31}$$

As discussed before, expression (29) has been obtained before by several authors by different methods. The results for the strong (σ_s) and interferency part (σ_{cs}) , eqs. (30) and (31), we believe to be novel.

From the qualitative point of view it is clear that the theoretical prediction for the $b_{min} = r_0(A_A^{\frac{1}{3}} + A_B^{\frac{1}{3}})$ dependence of the total strong cross section (eq.(30)) is in contradiction with the empirical parametrizations obtained in the E814 collaboration given by (Ref.[8]):

$$\sigma_{exp-s} = a(A_A^{1/3} + A_B^{1/3}) \tag{32}$$

where $a = 1.34 \pm 0.19 fm^2$. The theoretical cross sections depend very sensitively on the realistic values of the exchanged meson mass $\mu(MeV)$ and of the coupling constant $\frac{g_v^2}{(4\pi)} = 0.9$ [20]. For these reasons the nuclear and interference contributions to the collective excitation become negligible. Also as expected these two contributions are energy independent. The conclusion then is that the collective effects are essentially due to coulomb excitation and that the proposed mechanism is in rather good agreement with the data (see table 1, columns 2 and 4). Notice, however, that the experimental values for the strong part

target	σ_{exp-c}	σ_{exp-s}	σ_c	σ_{cs}	σ_s
C	6.81 ± 1.03	7.13 ± 1.00	6.34	1.13×10^{-2}	7.00×10^{-3}
Al	28.15 ± 4.88	8.08 ± 1.14	30.71	6.68×10^{-3}	1.70×10^{-3}
Cu	122.79 ± 24.16	9.40 ± 1.33	131.04	2.63×10^{-3}	1.69×10^{-4}
Sn	333.87 ± 70.92	10.64 ± 1.50	366.79	9.14×10^{-4}	1.51×10^{-5}
Pb	828.25 ± 187.65	12.00 ± 1.70	927.11	2.50×10^{-4}	9.01×10^{-7}

Table 1: *Experimental Coulomb (σ_{exp-c}), Experimental Strong (σ_{exp-s}), our Coulomb (σ_c, eq.29) ,Interference (σ_{cs}, eq.31), and our Strong cross section in the last column (σ_s, eq.30), with $g_v = 3.36$ and $\mu = 300 MeV$. Units in mb.*

of the interactions are non negligible. This indicates that one should rather investigate incoherent contributions to the nuclear part. In fact it can be shown that the incoherent knock-out process gives both qualitative and quantitative agreement with the empirical findings.

2.3   Conclusions for Coherent Effects

In this section we use a covariant formulation to describe the excitation of a collective mode via photon and meson exchange. Besides the well known total Coulomb contribution, analytical results are obtained for the total strong interaction contribution and the interference cross section. We show that the fact that the strong interaction is mediated by massive mesons drastically reduces its contribution to the excitation of the giant resonance. Both the pure strong and interference contributions are independent of the incident energy.

Moreover the theoretical prediction for the strong part of the collective total cross section is in qualitative disagreement with the experimental findings.

Therefore we conclude, that the electromagnetic effect is the only responsible for the processes where coherent effects are produced, as a giant dipole resonance. The strong interaction contribution is coming from incoherent effects such as one nucleon knock-out.

This point is investigated in the next section.

3   Nuclear Incoherent Effects

In this section we analyse Glauber's formulation of the one nucleon removal process according to reference [14], and give a theoretical expression for the total nuclear cross section. We show its compatibility with the empirical formula used in the analysis of the E814 experiment [8].

The physical process considered is a nucleon knock-out off the projectile within the context of Glauber's formalism.

3.1   The Nuclear One Nucleon Removal Cross Section

In this section the contribution of the strong interaction for the one nucleon removal process is calculated using Glauber's multiple scattering formalism [14]. The notation is summarized in the following equation, valid in the projectile's rest frame:

$$
^{A}Z\{|\vec{0}, \psi_0 >\} + T\{|-\vec{p}_0, \theta_0 >\} \quad \rightarrow^{A-1} Z\{|\vec{k}, \phi_\alpha >\}
$$
$$
+ n\{|\vec{p}, \eta_{\vec{p}} >\} + T'\{|-\vec{p}_0 - \vec{q}, \theta_\beta| >\} \qquad (33)
$$

Before the reaction, the projectile ^{A}Z, with intrinsic wave function ψ_0 is at rest and the target approaches with momentum $-\vec{p}_0$. After the interaction, a fragment with $A-1$ nucleons is detected in a particle stable state ϕ_α with momentum $\vec{k}$. The nucleon with momentum $\vec{p}$ and scattering wave function $\eta_{\vec{p}}$, and the final state Θ_β of the target, are unobserved. The cross-section corresponding to the equation (33) is given by [13]

$$\frac{d^3\sigma}{dk^3} = \int d^2q \sum_\alpha \sum_\beta$$

$$|\int \frac{d^2b}{2\pi} e^{-i\vec{q}\cdot\vec{b}} \left\langle \eta_{\vec{q}-\vec{k}}; \Theta_\beta | 1 - \prod_{i\in P, j\in T} (1-\Gamma_{ij}) |\psi_0; \Theta_0 \right\rangle|^2 \qquad (34)$$

In the above formula $\Gamma_{ij}(\vec{x}_i + b - \vec{y}_j)$ are the profile functions for collisions between a projectile nucleon i located at $\vec{x}_i$ and a target nucleon j at $\vec{y}_j$; $\vec{x}_i$ and $\vec{y}_j$ refer to the respective centers of mass and the impact parameter $\vec{b}$ is the transverse distance between projectile and target.

The fragmentation cross-section corresponding to the incoherent process is derived in detail in reference [14]. We quote the result

$$\frac{d^3\sigma}{dk^3} = \int D(\vec{s})d^2s \int dz \int d^3k W(\vec{s}, z, \vec{k}) \qquad (35)$$

where $\vec{s}$ is the perpendicular distance of each nucleon to the center of the nuclei. The distorting function $D(\vec{s})$ contains the reaction dynamics

$$D(\vec{s}) = \int d^2b \, e^{-2Im\chi_{FT}(\vec{b})-2Im\chi_T(\vec{b}+\vec{s})} (e^{\sigma_{NN}(0)T_T(\vec{b}+\vec{s})} - 1) \qquad (36)$$

where

$$i\chi_{FT}(\vec{b}) \approx -\langle \phi_\alpha \Theta_0 | \sum_{m\geq 1} \Gamma_{mn} | \phi_\alpha \Theta_0 \rangle \qquad (37)$$

$$i\chi_T(\vec{b}+\vec{s}) \approx -\langle \Theta_0 | \sum_n \Gamma_{1n}(\vec{b}+\vec{s}-\vec{t}_n) | \Theta_0 \rangle \qquad (38)$$

$$T_T(\vec{B}) = A_T \int dz \rho(\vec{B}, z) \qquad (39)$$

with $\rho(\vec{B}, z)$ is the nucleon density function, and

$$\sigma_{NN} = \int d^2q \, e^{i\vec{q}\cdot\vec{b}} \frac{d\sigma_{NN}}{d^2q} \qquad (40)$$

is the nucleon-nucleon cross section.

As discussed in reference ([14]) shown below, $D(\vec{s})$ localizes the reaction to the nuclear surface. The Wigner transform $W(\vec{s}, z, \vec{k})$ contains the momentum distribution of the removed nucleon.

In the calculations we have used Gaussian densities,

$$\rho(\vec{r}) = N_g e^{-(r^2/a_0^2)} \tag{41}$$

where for light nuclei we have adjusted a_0 in order to reproduce the root mean square radius. For heavier nuclei ($A > 40$) N and a_0 have been adjusted in order to reproduce the tail of the Woods-Saxon density.

The theoretical results to be shown are based on the result

$$\sigma_N = \int 2\pi s D(\vec{s}) T_p(\vec{s}) ds \tag{42}$$

with

$$T_p(\vec{s}) = \int dz \int_{k_{min}}^{\infty} d\vec{k} \, W(\vec{s}, z, \vec{k}) \tag{43}$$

k_{min} corresponds to the Coulomb barrier cut-off parameter, in the case where the nucleon removed is charged (σ_N is the nuclear cross section).

3.2 Results - E814 Brookhaven Experiment

In the E814 Experiment was reported the results of the EMD of ^{28}Si into $p + {}^{27}Al$, and Pb, Sn, Cu, Al and C were used as targets. The total cross section for producing $p + {}^{27}Al$ was measured. Presenting the first measurements of the final-state energy spectrum (for all targets) which is peaked near the isovector giant-dipole resonance in ^{28}Si. The incident energy was 14.5 GeV/nucleon. A detailed description of the experiment can be found in ref.[8]. The empirical fitting was for the electromagnetic part:

$$\sigma_{exp-c} = bZ_T^c \tag{44}$$

and for the nuclear part:

$$\sigma_{exp-s} = a(A_T^{1/3} + A_P^{1/3}) \tag{45}$$

with $a = 1.34 \pm 0.19 fm^2., b = 0.23 \pm 0.05 fm^2.$ and $c = 1.8 \pm 0.06..$ Using this fitting the authors showing a good agreement for all ions. We calculate the nuclear incoherent cross section (eq.42), using the corresponding cut-off for the Coulomb barrier.

Before a comparison between the theoretical results and the E814 data a word about the adequacy of such theoretical description to this specific set of data is in order: in the E814 experiment events which correspond to an energy transfer above 200 MeV have been excluded precisely in order to avoid contributions coming from the nuclear interaction. In Glauber's model the energy transfer is realized by means of nucleon-nucleon collisions which are assumed to occur essentially in the impact parameter plane. The average momentum transfer is 400 MeV/c which corresponds to and energy of about 80 MeV. Also the highest probabilities occur for smaller momentum transfers. We can therefore not exclude such processes from the experiment. The qualitative and quantitative agreement with experiment is, in the end what strengthened our conclusion.

In Table 2 we can see the nuclear cross section fitted by the experimentalists and our theoretical calculation of the nuclear part. We find good agreement. In this way our conclusion is that the processes is due to incoherent effects. Our numerical result based in eq. (42) gives a nuclear cross section proportional to the sum of the radii of the target and projectile, this is showing in the third coulumn of this table where the ratio $\frac{\sigma_N}{(A_P^{1/3}+A_T^{1/3})}$ is shown to remain constant.

We demonstrated that in this class of process, the interference cross section between Coulomb and strong interaction is negligible. The nuclear part is due to an incoherent process. In the context of Glauber's theory, we show that the nuclear cross section is proportional to the sum of the radii of nuclei, in perfect agreement with the experimental fitting used in E814 Collaboration.

In table 3 we can see the total cross section obtained in the E814 Collaboration in comparison with our theoretical results for the electromagnetic coherent and nuclear in-

A_T	$\sigma_{exp-s}(E814)$	σ_N	$\sigma_N/(A_P^{1/3} + A_T^{1/3})$
^{12}C	7.11 ± 1.00	8.11	1.5
^{27}Al	8.04 ± 1.14	9.24	1.5
^{63}Cu	9.34 ± 1.33	11.12	1.5
^{118}Sn	10.59 ± 1.50	13.04	1.5
^{208}Pb	11.95 ± 1.70	15.23	1.6

Table 2: *First column shows the parametrization used in E814 (nuclear part), second column our results, third column shows the proportionality with the sum of the radii.*

A_T	$\sigma_{exp-t}(E814)$	$\sigma_N + \sigma_{cs}$
^{12}C	13.92 ± 2.03	14.45
^{27}Al	36.19 ± 6.02	39.95
^{63}Cu	132.13 ± 25.49	142.16
^{118}Sn	344.46 ± 72.47	349.83
^{208}Pb	840.20 ± 189.35	942.34

Table 3: *First column shows the total cross section data of the E814 Collaboration, second column show the sum of our results for the nuclear and electromagnetic cross section.*

coherent parts. Good agreement is found.

4   Conclusions

In the present work the nuclear and interference nuclear-coulomb contribution to one nucleon removal processes in peripheral RHI collisions are studied. Firstly its coherent contribution is analyzed assuming that a giant dipole resonance is excited via one vector meson exchange. To lowest order simple analytical expression for the electromagnetic, nuclear and coulomb-nuclear interference contributions are derived.

For the electromagnetic part a cross section proportional to $\approx Z^2$ is obtained, as is expected experimentally and theoretically. The interference and coherent nuclear part

are small and in qualitative disagreement with the data. We therefore conclude that the nuclear contribution in experiment of one nucleon dissociation, results from incoherent nucleon-nucleon collisions. Therefore, in section 3, the incoherent one nucleon knock-out is studied within the context of Glauber's formalism. The total nuclear contribution so obtained is shown to be proportional to the sum of the radii of the nuclei as could be intuitively expected. Quantitatively speaking the incoherent process is shown to be of the order expected in experiments.

5   Acknowledgements

References

[1] H.H. Heckman e P. J. Lindstrom, Phys. Rev. Lett. **37** (1976) 56.

[2] D. L. Olson, B.L.Berman, D. E. Greiner, H. H. Heckman, P. J. Lindstrom, G. D. Westfall and J. J. Crawford, Phys. Rev.**C24** (1981) 1529.

[3] M. T. Mercier, Jhon C. Hill e F. K. Wohn, Phys. Rev. Lett. **52** (1984) 898.

[4] John C. Hill, F. K. Wohn and D. D. Schwellenbach, Proc. 2nd. International Workshop on Relativistic Aspects of Nuclear Physics, Brazil,1991, ed. T.Kodama (World Scientific, 1991) p.255

[5] M. T. Mercier, John C. Hill, F. K. Wohn, C. M. McCullough, M. E. Nieland, J. A. Winger, c. B. Howard, S. Renwick, D. K. Matheis and A. R. Smith, Phys. Rev. C **33** (1986) 1655.

[6] John C. Hill, F. K. Wohn, J. A. Wigner, M. Khayat, M. T. Mercier and A. R. Smith, Phys. Rev. C **39** (1989) 524.

[7] John C. Hill, F. K. Wohn, J. A. Winger and A. R. Smith, Phy. Rev. Lett. **60** (1988) 999.

[8] J. Barrette, R. Bellwied, P. Braun-Munzinger, W.E. Cleland, G. David, J. Dee, O. Dietzsch, E. Duek, M. Fatyga, D. Fox, S.V. Greene, J.R. Hall, T.K. Hemmick, N. Herrmann, R.W. Hogue, B. Hong, K. Jayananda, D. Kraus, B. Shiva Kumar, R. Lacasse, D. Lissauer, W.J. Llope, T. Ludlam, R. Majka, D. Makowiecki, S.K. Mark, S. McCorkle, J.T. Mitchell, D. Muthuswamy, E. O'Brien, V. Polychronakos, C. Pruneau, F.S. Rotondo, J. Sandweiss, J. Simon-Gillo, U. Sonnadara, J. Stachel, H. Takai, E.M. Takagui, T.G. Throwe, L. Waters, W.J. Willis, C. Winter, K. Wolf, D. Wolfe, C.L. Woody, N. Xu, Y. Zhang, Z. Zhang and Z. Zou, Phys. Rev. C **45** (1992) 2427.

[9] J. Barrette, P. Braun-Munzinger, W.E. Cleland, G. David, E. Duek, M. Fatyga, D. Fox, S.V. Greene, J.R. Hall, R. Heifetz, T.K. Hemmick, N. Herrmann, R.W. Hogue, G. Ingold, K. Jayananda, D. Kraus, A. Legault, D. Lissauer, W.J. Llope, T. Ludlam, R. Majka, D. Makowiecki, S.K. Mark, J.T. Mitchell, D. Muthuswamy, E. O'Brien, L.H. Olsen, V. Polychronakos, M. Rawool-Sullivan, F.S. Rotondo, J. Sandweiss, B. Shiva-Kumar, J. Simon, U. Sonnadara, J. Stachel, J. Sunier, H. Takai, T.G. Throwe, H. VanHecke, L. Waters, W.J. Willis, K. Wolf, D. Wolfe and C.L. Woody, Phys. Rev. C **41** (1990) 1512.

[10] L. Gonzaga Ferreira Filho, T. Kodama and M. C. Nemes, Nucl. Phys. **A 499** (1989) 837.

[11] R. J. Glauber, in *Lectures in Theoretical Physics*, edited by W. E. Brittin and L. G. Dunham (Interscience, New York, 1959), Vol. 1.

[12] R. J. Glauber and G. Matthiae, Nucl. Phys. B **21** (1970) 135.

[13] T. Fujita and J. Hufner, Nucl. Phys A **343** (1980) 493.

[14] J. Hufner and M.C. Nemes, Phys. Rev. C **23** (1981) 2538.

[15] C. A. Bertulani and G. Bauer, Phys. Rep. (1988) 163. 299.

[16] D. Galetti, T. Kodama and M. C. Nemes, Ann. Of Phys. **177** (1987) 229.

[17] C.J. Benesh, B.C. Cook and J.P. Vary, Phys. Rev. C, **40** (1989) 1198.

[18] A. Eiras and M. C. Nemes, J. Phys. G**20** (1994) 261.

[19] T Suzuki and D. J. Rowe, Nucl. Phys. **A286** (1977) 307.

[20] R. Machleidt, K. Holinde and Ch. Elster. Phys. Rep. 149 (1987) 1.